

# On the two kinds of vector particles

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All known elementary vector particles, the photon,  $Z$ ,  $W$  and the gluons, are described by the gauge theory. They belong to the *real* representation  $(1/2, 1/2)$  of the Lorentz group. On the other hand *inequivalent* representations  $(1, 0)$  and  $(0, 1)$  also correspond to particles with spin 1. It is natural to suppose that, along with the known vector particles, the new particles can exist. Evidence for the existence of these particles in nature is the presence of the axial-vector meson resonances with quantum numbers  $1^{+-}$ . Other indications for their existence are discussed. The signatures of their contributions into different physical processes are presented.

## I. INTRODUCTION

The description of the elementary particles relies on concepts of symmetry [1]. The rotation invariance group gives a particle classification with respect to the spin. In the quantum theory the spin may have half-integer or integer value. The lowest representation of the  $O(3)$  group, which can be used as a building block for the construction of a higher spin representation, is a two-component spinor  $\psi_\alpha$  ( $\alpha = 1, 2$ ). It describes the particles with spin  $1/2$ . In non-relativistic quantum mechanics there exists only one possibility to construct a vector:  $\mathbf{1}/2 + \mathbf{1}/2 = \mathbf{1}$ .

In the relativistic theory the symmetry group is the Lorentz group  $O(3, 1)$ , which is isomorphic to the direct product of the two spatial rotation groups  $O(3) \times O(3)$ . Therefore, two inequivalent representations  $(1/2, 0)$  and  $(0, 1/2)$  exist for the spin  $1/2$ . They correspond to particles with different chiralities and are represented by the Weyl spinor  $\psi_\alpha$  and its conjugate  $\psi_\alpha^* \equiv \psi_{\dot{\alpha}}$ , which are related by the  $P$  transformation. There exist two possibilities to construct the representation of the spin 1.

The vector representation  $(1/2, 1/2)$  is the well studied one and corresponds to the gauge particles. It is chirally neutral, because it arises from the product between the left  $(1/2, 0)$  and the right  $(0, 1/2)$  fundamental spinors. This property reflects the simple fact that all gauge interactions preserve chiralities of incoming and outgoing particles. This representation is transformed as a mixed (dotted and undotted) spinor  $\phi_{\alpha\dot{\beta}}$ , which is equivalent to the Lorentz vector

$$V_m = (\sigma_m)^{\alpha\dot{\beta}} \phi_{\beta\dot{\alpha}}, \quad (1)$$

where  $\sigma_i$  ( $i = 1, 2, 3$ ) are the Pauli matrices and  $\sigma_0$  is the unit matrix.

Let us stress that *inequivalent chiral* representations  $(1, 0)$  and  $(0, 1)$ , which also correspond to particles with spin 1, can be constructed if one uses only the product either of the left  $(1/2, 0)$  or of the right  $(0, 1/2)$  fundamental spinors. They are transformed independently by the proper Lorentz group as rank-2 spinors, symmetric in the spinor indices:  $\phi_{\alpha\beta}$  and  $\phi_{\dot{\alpha}\dot{\beta}}$ , respectively. To pass to more convenient Lorentz indices, the decomposition of the product of the Pauli matrices into symmetric and antisymmetric parts can be used:

$$(C\hat{\sigma}_m\sigma_n)^{\alpha\beta} = g_{mn}C^{\alpha\beta} - \frac{i}{2}\epsilon_{mnab}(C\hat{\sigma}^a\sigma^b)^{\alpha\beta}, \quad (2a)$$

$$(\sigma_m\hat{\sigma}_n C)^{\dot{\alpha}\dot{\beta}} = g_{mn}C^{\dot{\alpha}\dot{\beta}} + \frac{i}{2}\epsilon_{mnab}(\sigma^a\hat{\sigma}^b C)^{\dot{\alpha}\dot{\beta}}, \quad (2b)$$

where  $(\hat{\sigma}_m)_{\alpha\dot{\beta}} = (C^{-1}\sigma_m^T C)_{\alpha\dot{\beta}}$ ,  $C_{\alpha\beta} \equiv \epsilon_{\alpha\beta}$  is the charge-conjugation matrix, and  $\epsilon^{mnab}$  is the completely antisymmetric tensor, with  $\epsilon^{0123} = +1$ . Therefore, one can introduce the antisymmetric anti-self-dual tensor  $T_{mn}^- = T_{mn} - \tilde{T}_{mn}$ :

$$T_{mn}^- = \frac{i}{2}\epsilon_{mnab}(C\hat{\sigma}^a\sigma^b)^{\alpha\beta}\phi_{\beta\alpha}, \quad (3a)$$

and the antisymmetric self-dual tensor  $T_{mn}^+ = (T_{mn}^-)^* = T_{mn} + \tilde{T}_{mn}$ :

$$T_{mn}^+ = -\frac{i}{2}\epsilon_{mnab}(\sigma^a\hat{\sigma}^b C)^{\dot{\alpha}\dot{\beta}}\phi_{\dot{\beta}\dot{\alpha}}, \quad (3b)$$

where  $\tilde{T}_{mn} = (i/2)\epsilon_{mnab}T^{ab}$  is the tensor that is dual to the real antisymmetric tensor  $T_{mn}$ . Hence, the real tensor  $T_{mn}$  corresponds both to the left and to the right vector particles or, in other words, to the vector and axial-vector ones.

As far as the vector potential  $V_m$  (1) describes the real vector particles, the antisymmetric tensors  $T_{mn}^\pm$  (3)

can correspond to yet unknown vector particles. It is clear that the transformation law of the vector particles uniquely defines their trilinear renormalizable interaction with the fermions without derivatives. For example, the usual gauge interaction term  $\bar{\Psi}\gamma^m\Psi V_m$  arises from the Lorentz-invariant form  $\psi^\alpha\phi_{\alpha\dot{\beta}}\psi^{\dot{\beta}}$ , where

$$\Psi = \begin{pmatrix} \psi_\alpha \\ iC^{\dot{\alpha}\beta}\psi_{\dot{\beta}} \end{pmatrix} = \begin{pmatrix} \Psi_L \\ \Psi_R \end{pmatrix}, \quad \gamma^m = \begin{pmatrix} 0 & \hat{\sigma}^m \\ \sigma^m & 0 \end{pmatrix},$$

are the Dirac bispinor and the Dirac matrices in the helicity representation. On the other hand, the trilinear interaction of the new vector particles with the fermions  $\psi^\alpha\phi_{\alpha\beta}\psi^\beta + \psi^{\dot{\alpha}}\phi_{\dot{\alpha}\dot{\beta}}\psi^{\dot{\beta}}$  corresponds to the Yukawa term\*  $\bar{\Psi}\sigma^{mn}\Psi T_{mn}$ , where  $\sigma^{mn} = i[\gamma^m, \gamma^n]/2$ . The key feature of the interactions mediated by the new vector particles is chirality flip of incoming and outgoing particles. The gauge vector particles and the new ones have different interactions and, consequently, different quantum numbers. Therefore, they correspond to different kinds of vector particles. A common opinion exists in the literature, that massive vector particles can be equivalently described by the vector potential  $V_m$  or by the antisymmetric tensor field  $T_{mn}$  [2]. This is true for free particles as far as both representations correspond to particles with spin 1. However, there is no equivalence when an interaction is included. Only interactions, in particular with fermions, can distinguish between these two kinds of vector particles. To reveal this let us consider the hadron vector resonances.

## II. HADRON VECTOR RESONANCES

The bound states of a quark and an antiquark are characterized by the quantum numbers  $J^{PC}$ , where  $J$  is the total angular momentum,  $P$  is the parity and  $C$  is the charge conjugation. There exist three types of different quantum numbers for the known vector mesons [3]. They are  $1^{--}$ ,  $1^{++}$  and  $1^{+-}$ . For instance, the first quantum numbers are assigned to the  $\rho$ ,  $\omega$  and  $\phi$  vector mesons. The second quantum numbers are assigned to the  $a_1$  and  $f_1$  axial-vector mesons; however, note that the third quantum numbers are assigned again to the axial-vector mesons  $b_1$  and  $h_1$ . The key point is the difference between the last two assignments for the axial-vector mesons.

Let us consider the extended Nambu–Jona-Lasinio (ENJL) [4] models. In such models the Lagrangian contains only the quark fields, while the spontaneous sym-

metry breaking and the hadron states are generated dynamically by the model itself. The mesonic states arise as excitations of quark–antiquark pairs and that defines their interactions with the quarks. There are the vector  $\bar{\Psi}\gamma^m\Psi$  and the axial-vector  $\bar{\Psi}\gamma^m\gamma^5\Psi$  bilinear forms of the quark spinor fields with quantum numbers  $1^{--}$  and  $1^{++}$ , which correspond to the vector and axial-vector mesons, respectively. They have gauge-like interactions with the quarks and can be described by the gauge vector  $V_m$  and axial-vector  $A_m$  fields. Up to now all local ENJL models include only vector mesons with quantum numbers  $1^{--}$  and  $1^{++}$  and do not describe mesons with quantum numbers  $1^{+-}$ . One way of incorporating these mesons into the ENJL model has been described in [5]. These mesonic states correspond to the vector particles, which are described by the antisymmetric tensor field  $T_{mn}$ , rather than by the gauge fields. The bilinear form  $\bar{\Psi}\sigma^{mn}\Psi$  is used to describe the quantum numbers and the interactions of these mesons. The existence of the axial-vector mesons with the quantum numbers  $1^{+-}$  points out that the new kind of vector particles should exist in nature.

What concerns the hadron physics the introduction of the antisymmetric tensor field in the ENJL model can give new understanding of the vector mesonic resonances. The six components of the antisymmetric tensor correspond to the axial-vector  $B_m = \partial_n(\bar{\Psi}\sigma^{mn}\gamma^5\Psi)$  with quantum numbers  $1^{+-}$  and also to the vector  $R_m = \partial_n(\bar{\Psi}\sigma^{mn}\Psi)$  with quantum numbers  $1^{--}$ . Each of these vectors  $B_m$  and  $R_m$  has three independent components due to the antisymmetric property of  $\sigma^{mn}$ . Therefore, besides the axial-vector mesons  $B_m$  with quantum numbers  $1^{+-}$ , there are additional vector mesons  $R_m$  with quantum numbers  $1^{--}$ , like those of the gauge mesons  $V_m$ , but having different coupling to the quarks. As far as there exist two different vector particles with the same quantum numbers  $1^{--}$ , their mass eigenstates can be a linear combination of them<sup>†</sup>. For example, for the isospin 1 vector mesons they are  $\rho(770)$  and  $\rho'$ , where the latter is either the  $\rho(1450)$  or  $\rho(1700)$  state. For the axial-vector mesons it is possible to make a unique identification of the low-lying mesonic states with the quantum numbers  $1^{++}$  and  $1^{+-}$  as  $a_1(1260)$  and  $b_1(1235)$ , respectively. Applying the ENJL approach to all these (axial)-vector mesons, a remarkable relation among their masses is derived [5]:

$$m_\rho^2 + m_{\rho'}^2 = m_{a_1}^2 + m_{b_1}^2. \quad (4)$$

The extraction of the  $a_1$  and  $\rho'$  masses from the experiments is a model-dependent procedure. The masses have a broad range of values. Therefore, to fix one of these

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\*For the gauge antisymmetric tensor field  $B_{mn}$  the gauge-invariant form of the interaction with derivative  $\epsilon^{abcm}\partial_{[a}B_{bc]}\bar{\Psi}\gamma_m\Psi$  is used, which is non-renormalizable. Here we will consider the antisymmetric tensor matter field  $T_{mn}$  with a Yukawa interaction.

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<sup>†</sup>It means, in particular, that the  $\rho$  meson can have both gauge and anomalous tensor couplings with the quarks, while the axial-vector mesons with quantum numbers  $1^{++}$  have only gauge interactions and the axial-vector mesons with quantum numbers  $1^{+-}$  have only tensor interactions.

masses one can use as approximation the Weinberg sum relation [6]

$$m_{a_1} \simeq \sqrt{2} m_\rho. \quad (5)$$

This fixes the  $a_1$  mass value  $m_{a_1} \simeq 1088$  MeV, and favours the choice of  $\rho(1450)$  resonance as the second mass eigenstate with the quantum numbers  $1^{--}$ , because eq. (4) gives  $m_{\rho'} \simeq 1450$  MeV. Certainly, the new point of view on the hadron (axial)-vector resonances enables us not only to explain their mass spectrum, but also introduces a new decay dynamics connecting with their tensor couplings to the quarks.

### III. ELECTROWEAK PHYSICS

Since two types of vector particles are used for the hadron resonance description, one can extrapolate this feature to the electroweak physics. Let us introduce a new type of intermediate vector bosons into the Standard Model (SM). As far as the tensor interaction does not conserve chirality, these bosons should couple to the left doublets and to the right singlets of the fermion fields. In order to have  $SU(2) \times U(1)$  invariance, the new vector bosons must be doublets. Therefore, they do not mix with the gauge bosons before the symmetry breaking. It means that in the high energy experiments they do not interfere with the gauge bosons and contribute to processes as their squared amplitude.

Nevertheless, due to the particular chiral properties of the new particles, the angular distribution of the fermion annihilation cross-section for the tensor coupling differs from the one of the vector coupling. For the vector case at high energies the well-known expression holds<sup>‡</sup>

$$\frac{d\sigma}{d\cos\theta} \sim (1 + \cos^2\theta). \quad (6a)$$

While the anomalous tensor interaction of the currents  $\partial_n(\bar{\Psi}\sigma^{mn}\Psi)$  leads to a different angular dependence

$$\frac{d\sigma}{d\cos\theta} \sim \cos^2\theta. \quad (6b)$$

The differential cross-section (6b) also differs from the case of a scalar coupling, which has uniform distribution.

This means that if the new kind of vector particles exists, at high energies they should give an essential contribution to the differential cross-section only near to the beam direction. It is not easy to detect definitely such an excess experimentally, but it will be a clear signature of the new kind of interactions. Some indications of such

excesses exist in recent experimental data with high pseudorapidity  $\eta \geq 2$ : in the production of the  $b$ -quarks at the Tevatron [7] and the  $\tau\bar{\tau}$  pairs production at LEP 2 [8].

The LEP experiments at CERN are very suitable for a search of such deviations from the SM. Generally, the lepton physics is free from uncertainties connected with parton distributions, which are typical for the hadron colliders. As far as the new interactions do not conserve chirality, their effects increase with the mass of the particles. Therefore, the analysis of heavy particle decays, such a  $\tau$  leptons, gives a unique possibility for searching these effects. It was noted in [9] that the new interactions lead to another parametrization of pure leptonic decays, different from the conventional Michel parametrization. The experimental analysis of the  $\tau$  decays and constraints on the new tensor coupling are presented in [10].

The low energy physics is able either to constrain or to indicate the presence of the new kind intermediate particles. This can be realized mainly in precision experiments of particle decays. However, at present they do not obtain the deserved attention. It is noteworthy to point out here that, for example, the principal  $\rho$  parameter in the muon decay has not been measured since 1969. Needless to say that a knowledge of the precision spectrum shape is important for the determination of the Fermi coupling constant  $G_F$ .

If some new interactions give a contribution into the muon decay, this will effectively lead to a lower value of  $G_F$  in comparison with the experimentally extracted one, which assumes only SM interactions. It may help to solve the long-standing problem of a violation of unitarity in the first row of the Cabibbo–Kobayashi–Maskawa (CKM) matrix [11]:

$$|V_{ud}|^2 + |V_{us}|^2 + |V_{ub}|^2 = 0.9968 \pm 0.0014. \quad (7)$$

The main contribution into the unitarity sum (7) comes from the  $V_{ud}$  element of the CKM matrix. This element is extracted with high precision from the nuclear superallowed beta decays  $0^+ \rightarrow 0^+$ , comparing the strength of this vector transition to  $G_F$ :

$$|V_{ud}|^2 = \text{const}/G_F^2. \quad (8)$$

A two per mille lowered value of  $G_F$  will restore the unitarity condition.

In principle, the new tensor interactions can be responsible for such a scenario [9]. In this case the strength of these interactions must be comparable with the electroweak one and that can lead to additional observable contributions in other experiments. Let us point out some indication for a possible admixture of such interactions in radiative pion decay [12] and semileptonic three-body kaon decay [13].

First of all it is important to note that the tensor interactions do not contribute directly to the pion decay  $\pi \rightarrow e\nu$  and avoid the strong constraints on the pseudoscalar interactions. Although the tensor interactions,

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<sup>‡</sup>For simplicity the parity violation and the interference between the photon and  $Z$  in the SM can be neglected.

like the scalar ones, do not conserve chirality and should give a large contribution into chirality-suppressed pion decay, this is impossible because a tensor form factor for the pion matrix element  $\langle 0 | \bar{q} \sigma^{mn} q | \pi(p) \rangle$  cannot be constructed. A different situation arises when the radiative pion decay  $\pi \rightarrow e \nu \gamma$  is considered. In this case the tensor interactions may give a contribution themselves and, moreover, may interfere with a QED inner bremsstrahlung contribution. Exactly such a type of destructive interference was observed in the radiative pion decays in flight [12].

The other experiment, where the tensor form factor can be introduced on the same footing as the vector form factor of the SM, is the semileptonic three-body kaon decay  $K \rightarrow \pi e \nu$ . As far as most of the experiments on kaon decays give information only about the vector form factors, the experimental data are poor. However, the last high-statistics experiment [13] indicates the presence of non-zero tensor form factor.

In ref. [14] it was shown that both the anomalous destructive interference in the radiative pion decay and the presence of a non-zero tensor form factor in the semileptonic kaon decay, may be explained by an admixture of tensor interactions. This is also compatible with the contribution of the tensor interactions into the pure leptonic decays at the per mille level [9]. It is clear that new experiments are necessary to confirm or reject the presence of this new type of vector particles in nature.

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- [9] M. Chizhov, Mod. Phys. Lett. A9 (1994) 2979; On the muon decay parameters, hep-ph/9612399.
- [10] DELPHI Collaboration, A study of the Lorentz structure in tau decays, CERN-EP/2000-027, to be published in Eur. Phys. J. C.
- [11] W. Jaus and G. Rasche, Phys. Rev. D41 (1990) 166; I. S. Towner and J. C. Hardy, The current status of  $V_{ud}$ , nucl-th/9809087.
- [12] V. N. Bolotov et al., Phys. Lett. B243 (1990) 308.
- [13] S. A. Akimenko et al., Phys. Lett. B259 (1991) 225.
- [14] M. V. Chizhov, Mod. Phys. Lett. A8 (1993) 2753.

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- [1] Yu. V. Novozhilov, Introduction to Elementary Particle Theory, Pergamon Press, 1975
  - [2] J. Gasser and H. Leutwyler, Ann. Phys. 158 (1984) 142; G. Ecker et al., Nucl. Phys. B321 (1989) 311; Phys. Lett. B223 (1989) 425; D. Kalafatis, Phys. Lett. B313 (1993) 115.
  - [3] Particle Data Group, Eur. Phys. J. C15 (2000) 1.
  - [4] Y. Nambu and G. Jona-Lasinio, Phys. Rev. 122 (1961) 345.
  - [5] M. V. Chizhov, Tensor excitations in Nambu–Jona-Lasinio model, hep-ph/9610220.
  - [6] S. Weinberg, Phys. Rev. Lett. 18 (1967) 507.
  - [7] D0 Collaboration (B. Abbott et al.), Phys. Rev. Lett. 84 (2000) 5478.
  - [8] P. Renton, talk at the LEP Experiments Committee Meeting, July 2000.